A review on axions and the strong CP problem

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Abstract. We review the recent developments on axion physics. Some new comments on the strong CP problem, the axion mass, and the simple energy loss mechanism of white dwarfs and related issues are included.

INTRODUCTION

The most awaited information in the universe now is what is the DM of the universe. One plausible candidate is the weakly interacting massive particle(WIMP) and the other attractive candidate is a very light axion. We try to discuss these aspects from scarce experimental hints and comment a viable model in a SUSY framework. Axion is a Goldstone boson arising when the PQ global symmetry is spontaneously broken. The simple form dictates that its interaction is only through the gluon anomaly term $G\tilde{G}$. The axion models have the spontaneous symmetry breaking scale F and the axion decay constant F_a which are related by $F = N_{\mathrm{DW}}F_a$.

The WIMP was first discussed by Lee and Weinberg [1] where a heavy neutrino was considered, which was the beginning of the usage "weak" in WIMP. The LSP interaction is "weak" if interaction mediators the SUSY particles are in the 100 GeV range as W boson. That is the reason we talk about WIMP in TeV SUSY. At present, WIMP almost means the LSP.

On the other hand, if there exists a coherently oscillating boson field in the universe, this bosonic collective motion is always equivalent to CDM [2]. The dashed lines in Fig. 1(a) represent the shapes of CDMs of the WIMP and the bosonic collective motion.

Regarding the coherently oscillating bosonic field, the mostly discussed candidate is the very long lived axion [5, 6], suggested for the solution of the strong CP problem. The strong CP problem is that QCD allows a CP violating flavor singlet interaction which however is contradicted by the stringent bound [7] on the electric dipole moment of neutron(EDMn). The existence of instanton solution in nonabelian gauge theories needs θ vacuum [8]. It introduces an effective interaction, the so-called θ term,

$$\frac{\theta}{32\pi^2}G^a_{\mu\nu}\tilde{G}^{a\mu\nu} \tag{1}$$

where θ is the final value taking into account the electroweak CP violation and $G^a_{\mu\nu}$ is the field strength of gluon. For QCD to become a correct theory, this CP violating term must be sufficiently suppressed. Recently, taking the researches of last 20 years, the strong CP problem is reviewed in Ref. [9]. Our interest here is the bound on the
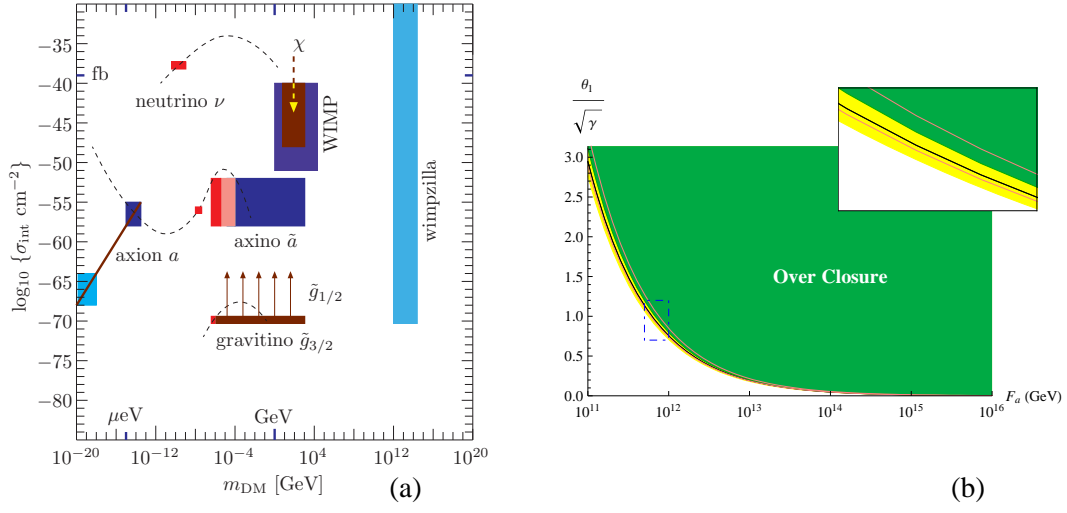
FIGURE 1. In (a), a rough sketch of masses and cross sections modified from [3] are shown. In (b), the axion among these is shown in the F_a versus the initial misalignment angle plane [4].

axion decay constant

$$10^9 \text{ GeV} \leq F_a \leq 10^{12} \text{ GeV}. \tag{2}$$

But the bound is in fact a two-dimensional region as shown in Fig. 1(b) [4], taking into account the anharmonic term carefully and the new mass bounds on light quark masses. Related studies can be found in [10]. Fig. 1(b) is the basis of using the anthropic argument for a large F_a [11].

STRONG CP PROBLEM AND SOLUTIONS

The so-called θ term resulting from the instanton background led to a sizable EDMn. But the observed EDMn is very tightly bounded, $|d_n/e| < 2.9 \times 10^{-26}$ cm [7]. In Ref. [9], a resulting θ bound has been given from an effective Lagrangian for the EDMn together with the magnetic dipole moment of neutron(MDMn). The fact that the mass term and the MDMn term have the same chiral transformation property has been used. The new bound is $|\theta| < 0.7 \times 10^{-11}$ which is a factor smaller than earlier bounds [12]. The strong CP problem is given here, "Why is θ term so small?" The strong CP problem has been solved in three different categories: (1) Calculable θ, (2) Massless up quark possibility, and (3) Axion solution.

- Calculable θ: The calculable solutions were very popular in 1978 [13]. But now only the Nelson-Barr model [14] has been remaining, which introduces vectorlike heavy quarks at high energy scale. This model produces the Kobayashi-Maskawa type weak CP violation at the low energy standard model(SM). Still, at one loop the appearance of θ must be forbidden up to one loop.

- Massless up quark: Suppose that we chiral-transform a quark q by $e^{i\gamma_5\alpha}q$,

$$\int\left(-m\bar{q}q+\frac{\theta}{32\pi^2}G\tilde{G}\right)\rightarrow\int\left(-m\bar{q}e^{2i\gamma_5\alpha}q+\frac{\theta-2\alpha}{32\pi^2}G\tilde{G}\right) \tag{3}$$

  If $m=0$, it is equivalent to changing $\theta\rightarrow\theta-2\alpha$. Thus, there exists a shift symmetry of $\theta\rightarrow\theta-2\alpha$. Then, θ is not physical, and there is no strong CP problem. The physical problem is, "Is massless up quark phenomenologically viable?" [15]. The recent compilation of the light quark masses gives, $m_u=2.5\pm1$ and $m_d=5.1\pm1.5$ in units of MeV [16], even not considering the lattice calculation of $m_u\neq0$ [17]. This is convincing enough that $m_u=0$ is not a strong CP solution.
- The axion solution is given below.

AXIONS

Historically, Peccei and Quinn (PQ) introduced the so-called PQ symmetry to mimic the symmetry of a massless quark of (3), $\theta\rightarrow\theta-2\alpha$, in the full electroweak theory [18]. The PQ symmetry includes the transformations of two Higgs doublets H_u and H_d, coupling to up-type and down-type quarks, respectively,

$$q\rightarrow e^{i\gamma_5\alpha}q,\ \{H_u,H_d\}\rightarrow e^{i\beta}\{H_u,H_d\} \tag{4}$$

for a symbolic Lagrangian, $\mathscr{L}=(\bar{q}_LH_uu_R+\bar{q}_LH_dd_R+\text{h.c.})-V(H_u,H_d)+\theta\{F\tilde{F}\}$, which changes to $\mathscr{L}=(e^{i(\beta-\alpha)}\bar{q}_LH_uu_R+e^{i(\beta-\alpha)}\bar{q}_LH_dd_R+\text{h.c.})-V(H_u,H_d)+(\theta-2\alpha)\{F\tilde{F}\}$. Therefore, for $\beta=\alpha$ PQ achieved almost the same thing as the massless quark case. The Lagrangian is invariant under changing $\theta\rightarrow\theta-2\alpha$; thus it seems that θ is not physical, since it is a phase of the PQ transformation. But, θ is physical. At the Lagrangian level, there seems to be no strong CP problem. But the VEVs $\langle H_u\rangle$ and $\langle H_d\rangle$ break the PQ global symmetry and there results a Goldstone boson, axion a [19]. Since θ is made a field, the original $\cos\theta$ dependence becomes the potential of the axion a. If its potential is of the $-\cos\theta$ form, always $\theta=a/F_a$ can be chosen at 0. So the PQ solution of the strong CP problem is that the vacuum chooses [18, 20]

$$\theta=0. \tag{5}$$

A historical note is that the above Peccei-Quinn-Weinber-Wilczek(PQWW) axion is ruled out early in one year [21]. Two years later after many tries of calculable solutions, the PQ symmetry came back incorporating heavy quarks Q, using a singlet Higgs field [5],

$$\mathscr{L}=(\overline{Q}_LSQ_R+\text{h.c.})-V(S,H_u,H_d)+\theta\{F\tilde{F}\} \tag{6}$$

where the Higgs doublets are neutral under the PQ transformation. If they are not neutral, then it is not necessary to introduce heavy quarks but a nontrivial PQ transformation for the Higgs doublets [6]. In any case, the axion is the phase of the SM singlet S if the VEV of S is much above the electroweak scale. The couplings of S determine the axion interactions. Because it is a Goldstone boson, the couplings are of the derivative form except the anomaly term.
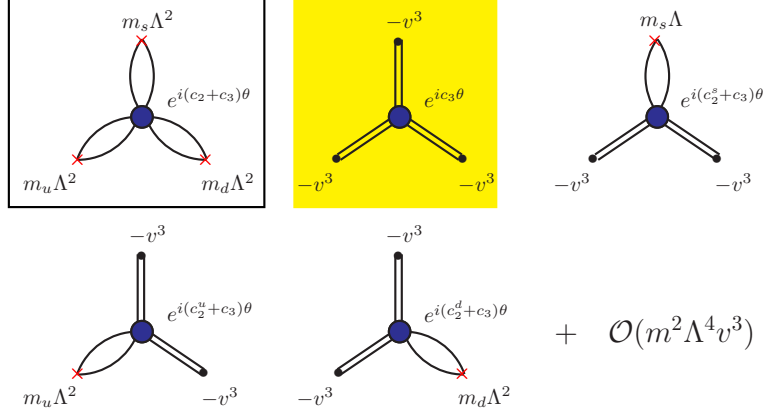
FIGURE 2. Contraction of quarks of the 't Hooft determinental interaction.

In most studies on axions, a specific example is chosen. Here, we consider a general description in an effective theory framework just above the QCD scale. Heavy fermions carrying color charges are special. All heavy fields, including Q and real S are integrated out. So consider the following Lagrangian,

$$
\begin{aligned}
\mathscr{L}_\theta &= \frac{1}{2} f_S^2 \partial^\mu \theta \partial_\mu \theta - \frac{1}{4g_c^2} G_{\mu\nu}^a G^{a\mu\nu} + (\bar{q}_L i \not{D} q_L + \bar{q}_R i \not{D} q_R) + c_1 (\partial_\mu \theta) \bar{q} \gamma^\mu \gamma_5 q \\
&\quad - \left(\bar{q}_L \, m \, q_R e^{ic_2\theta} + \text{h.c.} \right) + c_3 \frac{\theta}{32\pi^2} G_{\mu\nu}^a \tilde{G}^{a\mu\nu} \ (\text{or } \mathscr{L}_{\text{det}}) \\
&\quad + c_{\theta\gamma\gamma} \frac{\theta}{32\pi^2} F_{\text{em},\mu\nu}^i \tilde{F}_{\text{em}}^{i\mu\nu} + \mathscr{L}_{\text{leptons},\theta}
\end{aligned}
\tag{7}
$$

where $\theta = a/f_S$ with the axion decay constant f_S up to the domain wall number ($f_S = N_{DW} F_a$), q is the fermion matrix composed of $SU(3)_c$ charge carrying fields. There are three classes of couplings: c_1, c_2 and c_3. But the axion mass depends only on the combination of $c_2 + c_3$ [9]. The term $\mathscr{L}_{\text{det}}$ is the 't Hooft determinental interaction which is the basis of the solution of the U(1) problem. The determinental interaction is shown in Fig. 2 where the U(1) solution is given pictorially in the second diagram. If the story ends here, the axion is exactly massless. But, as shown in the other diagrams in Fig. 2, there are additional contributions which make axion massive. A 3×3 mass matrix of π^0, η', and axion can be diagonalized to give [9]

$$
m_{\pi^0}^2 = \frac{m_+ v^3 + 2\mu \Lambda_{\text{inst}}^3}{f_\pi^2}, \ m_{\eta'}^2 = \frac{4\Lambda_{\eta'}^4 + m_+ v^3 + 2\mu \Lambda_{\text{inst}}^3}{f_{\eta'}^2}, \ m_a^2 = \frac{Z}{(1+Z)^2} \frac{f_\pi^2 m_{\pi^0}^2 (1+\Delta)}{F_a^2}
\tag{8}
$$

where we used the parameters defined in [9]. The U(1) problem is solved by the first term of $m_{\eta'}^2$. In the axion mass, Δ represents the instanton contribution and $F_a = F/|c_2 + c_3|$ in terms of the singlet VEV scale F. Numerically, the axion mass in units of eV is $\simeq 0.6 \times 10^7 \text{ GeV}/F_a$. The essence of the axion solution is that $\langle a \rangle$ seeks $\theta = 0$ whatever happened before. In this sense it is a cosmological solution. The height of the potential is guessed as the QCD scale Λ_{QCD}^4.

TABLE 1. $c_{a\gamma\gamma}$ in several field theoretic models. (m,n) in the KSVZ denotes m copies of $Q_{\rm em} = \frac{2}{3}$ and n copies of $Q_{\rm em} = -\frac{1}{3}$ heavy quarks with the same PQ charge.

KSVZ *			DFSZ	
$Q_{\rm em}$	$c_{a\gamma\gamma}$	$x = \tan\beta$	same Higgs for (q^c, e) masses,	$c_{a\gamma\gamma}$
0	-1.95	any x,	(d^c, e)	0.72
$\pm\frac{1}{3}$	-1.28	any x,	(u^c, e)	-1.28

* For the unlikely cases of $Q_{\rm em} = \pm\frac{2}{3}, \pm 1$, we have $c_{a\gamma\gamma} = 0.72, 4.05$, respectively. For (m,m), we obtain $c_{a\gamma\gamma} = -0.28$.

Axion couplings

Above the electroweak scale, we integrate out heavy fields. If colored quarks are integrated out, its effect is appearing as the coefficient of the gluon anomaly. If only bosons are integrated out as in the DFSZ model, there is no such anomaly term. Thus, we have

$$
\begin{aligned}
&\text{KSVZ}: &&c_1 = 0,\ c_2 = 0,\ c_3 = \text{nonzero} \\
&\text{DFSZ}: &&c_1 = 0,\ c_2 = \text{nonzero},\ c_3 = 0 \\
&\text{PQWW}: &&\text{similar to DFSZ}
\end{aligned}
\tag{9}
$$

The axion-hadron couplings are important for the study of supernovae: The chiral symmetry breaking is properly taken into account, using the reparametrization invariance so that $c_3' = 0$ so that the axion-quark couplings are

$$
\text{KSVZ}: \qquad \bar{c}_1^{u,d} = \tfrac{1}{2}\bar{c}_2^{u,d},
\tag{10}
$$

$$
\text{DFSZ}: \quad \bar{c}_1^u = -\frac{v_d^2}{2(v_u^2+v_d^2)} + \tfrac{1}{2}\bar{c}_2^{u,d},\ \bar{c}_1^d = -\frac{v_u^2}{2(v_u^2+v_d^2)} + \tfrac{1}{2}\bar{c}_2^{u,d}
\tag{11}
$$

where $\bar{c}_2^u = 1/(1+Z), \bar{c}_2^d = Z/(1+Z)$ and $c_3' = 0$. The supernovae study of the KSVZ axion has been given before [22], and now the DFSZ axion can be studied also with the above couplings.

The next important axion coupling is the axion-photon-photon coupling which has a strong model dependence. For several different very light axion models, they are calculated in [26]. Some of these numbers are shown in Table 1.

Axion mixing in view of hidden sector

Even if we lowered some F_a from the GUT scale, we must consider the hidden sector also in SUSY models. In this case, the axion mixing must be considered. For the mixing, there is an important cross theorem on the decay constant: *Suppose two axions a_1 with F_1 and a_2 with F_2 ($F_1 \ll F_2$) couple to two nonabelian gauge groups whose scales have a hierarchy, $\Lambda_1 \ll \Lambda_2$. Then, diagonalization process chooses the larger condensation*
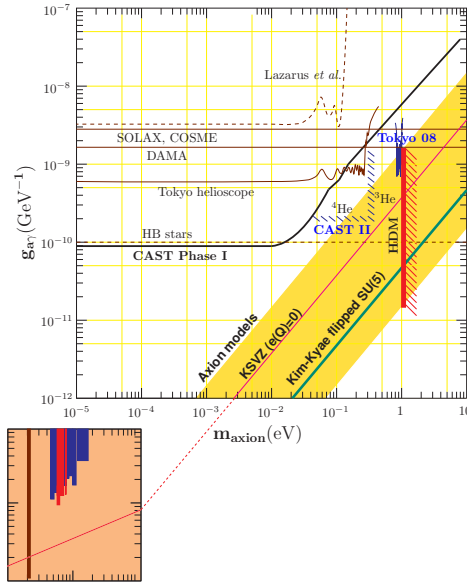
FIGURE 3. Astrophysical and cosmological bounds from axion experiments [9].

scale Λ_2 chooses the smaller decay constant F_1, and the smaller condensation scale Λ_1 chooses the larger decay constant F_2 [23]. So, just obtaining a small decay constant is not enough. The hidden sector may steal the smaller decay constant. It is likely that the QCD axion chooses the larger decay constant [24]. In this regard, we point out that the MI-axion with anomalous U(1) always has a large decay constant since all fields are charged under this anomalous U(1).

So, phenomenologically successful axion must need an approximate PQ symmetry. An approximate PQ global symmetry with a discrete symmetry in SUGRA was pointed out long time ago for Z_9 in [25]. But Z_9 is not possible in simple orbifold compactifications of string models. We may need a $Z_3 \times Z_3$ orbifold. From heterotic string, approximate PQ symmetry are considered in [27].

In Fig. 3, we show the current astrophysical and cosmological bounds on the axion decay constant. There exists one calculation [27] from a consistent string model containing a phenomenologically viable MSSM model [28], which is also shown in the figure.

White dwarf evolution

Among the star energy loss mechanisms by the axion emission, the recent white dwarf(WD) study is very interesting. WDs have a very simple structure and easy to study. Their chief aspect is a degenerate electron gas with the temperature $T/\rho^{2/3} < 1.3 \times 10^5 K \text{cm}^2 \text{gr}^{-2/3}$. For Sirius B, this number is 3.6×10^3 and it is a WD. The astronomers are able to recover the history of the star formation in our Galaxy by studying the statistics of WD temperatures. For this, the energy transport mechanism from the core of a WD is essential. Unlike in Sun, it is transported by neutrinos at high T since most electron are filling the degenerate energy levels. So, the transport mechanism

is very simple. And the resulting luminosity at the surface is calculable and reliable.

The later stage of the WD evolution is cristalization from the core. As time goes on, the luminosity drops. In terms of time t, the luminosity is $L_{\mathrm{WD}} = L_0(1 + \frac{t}{\tau_0})^{-7/5}$ where the characteristic time scale τ_0 can be taken as $\tau_0 \simeq 2.16 \times 10^7$ yrs, for example. A more complete treatment changes this simple behavior little bit [29]. The energy loss in the early stage of a WD is through the photon conversion to neutrino pairs in the electron plasma. This calculation of the photon decay was initiated in 1960s, but the accurate number was available after 1972 when the weak neutral current interaction was taken into account [30].

With more data, recently Isern *et al.* gives a very impressive figure from the most recent calculation of the above pioneering works, including the early stage and the crystalization period [31]. We translate their number to the axion-electron coupling strength for any axion model [32] and the resulting axion-electron Yukawa coupling,

$$\left| \frac{m_e \Gamma(e)}{F} \right| \simeq 0.7 \times 10^{-13}, \quad \frac{m_e \Gamma(e)/N_{\mathrm{DW}}}{F_a} \bar{e} i \gamma_5 e a \tag{12}$$

where $F = N_{\mathrm{DW}} F_a$ and $\Gamma(e)$ is the PQ charge of electron.

One can think of less exotic processes like the effects of the neutrino transition magnetic moments(NTMM, μ). For the NTMM and the weak neutral currents, one can consider the plasmon decay to neutrinos as [33],

$$\Gamma_{\mathrm{NTMM}} = \frac{\mu^2}{24\pi} Z_{T,L} \frac{(\omega_{T,L}^2 - \omega_{plasmon}^2)^2}{\omega_{T,L}} \tag{13}$$

$$\Gamma_{\mathrm{vector\ NC}} = \frac{G_F^2 C_V^2}{48\pi^2 \alpha_{\mathrm{em}}^2} Z_{T,L} \frac{(\omega_{T,L}^2 - \omega_{plasmon}^2)^3}{\omega_{T,L}}. \tag{14}$$

So, the ratio of radiation rates is expressed in terms of Q values (the convolution of the decay rates and the plasmon distribution function). Ref. [33] gives $Q_{\mathrm{NTMM}}/Q_{\mathrm{SM}} = O(1)$ for $\mu \sim 10^{-11}$ times the Bohr magneton. So, the problem here is of the extremely small NTMM ($\ll 10^{-11}\mu_{\mathrm{Bohr}}$) in the SM [34].

On the other hand, one may introduce new hypothetical light particles: (i) some kind of pseudo-Goldstone bosons considered by Haber [35], (ii) a massless or almost massless extra-photon with a kinetic mixing but without the electron coupling to the extra-photon, (iii) a sub-keV milli-charged particle, and (iv) a very light axion. For Case (i), the pseudo-Goldstone boson coupling to electron is required to be $(3-4) \times 10^{-13}$. For Case (ii), the extra-photon cannot be a candidate for the energy takeout of WDs. For Case (iii), the WD allowed parameter region of Ref. [36] is hardly achievable from the red giant bound [37]. Case (iv) is commented below.

To have a QCD axion at the axion window of $F_a \simeq 10^9 - 10^{12}$ GeV, we need some PQ charge carrying scalar developing VEV(s) at that scale. An enhanced electron coupling compared to the axion lower bound is possible by,

(*a*) Assigning a large PQ charge to *e*. But the quark-lepton unification makes this idea not very promising, especially in GUTs.
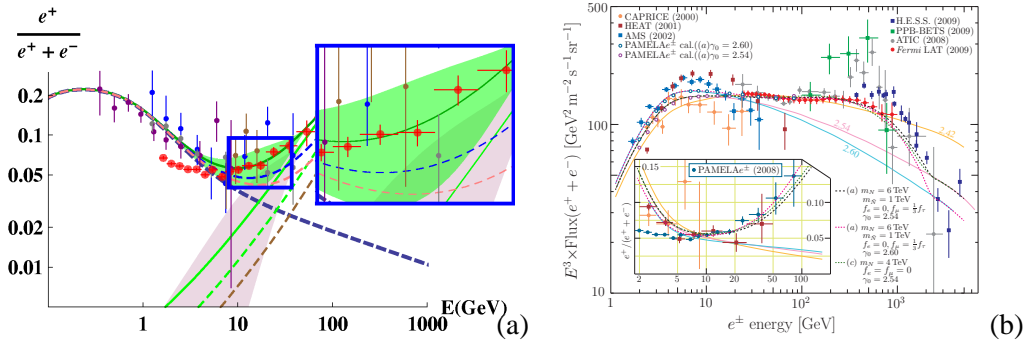
FIGURE 4. Positron excess near Earth: (a) scattering of two DMs, $\chi + N$, for the PAMELA e^+ [40] and (b) a decaying DM N [41] for the PAMELA e^+ and the *Fermi* LAT $e^+ + e^-$.

(*b*) Assign 1 PQ charge to e, but let the N_{DW} be fractional. In this case, only $\frac{1}{2}$ is possible. For the quark sector, effectively e.g. only one chirality of one quark carries the PQ charge. But both e_L and e_R carry the PQ charges.

Case (*b*) has been discussed in [32], where only u_R is the PQ charge carrying quark. This kind of model is possible in the flipped SU(5) since $(u, \nu_e, e)_L^T, u_R$ and e_R representations are independent. In addition, the flipped SU(5) model is good in that it is obtained from string compactification [38, 28] and also introduces e_R as a GUT singlet, making it possible to interpret the recent leptophilic PAMELA data. This sideway comment is that this flipped SU(5) provides a two DM scenario for the PAMELA positrons [39] through our singlet e_R which is promoted to a heavy charged lepton E. Ref. [40] was the earliest scattering model with a leptophilic property while saving neutralino χ as a DM component. Because the PAMELA data is so intriguing, here we show this idea very briefly from the flipped SU(5). Here, DMs are neutralino χ and a neutral chiral fermion N with the superpotential $W = f N_R E_R^c e_R^c + N_R^3$. In this case, a large enhancement factor is not needed. In the same model without the N_R^3 term and raising the mass of E to a GUT scale, the possibility of decaying DM N has been studied also [41] in view of the PAMELA and the *Fermi* LAT data [42]. For a decaying DM N, its number density can be calculated in models with a very heavy axino [43]. Since the WD axions and the axino are present in models of the MSSM extension with the PQ symmetry, these are all in the same framework. In Fig. 4, we show the allowed parameter regions of these studies.

Axions in the universe

The approximate form of the axion potential is of the cosine function, $-\cos\theta$. The minimum is at $\theta = 0$. But the axion potential is almost extremely flat and hence the vacuum stays at any θ for a long time. It starts to oscillate when the Hubble time $1/H$ is larger than the oscillation period of the classical axion field $\langle a \rangle$, $3H < m_a$. This occurs when the temperature is about 0.92 GeV [4]. In other words, the axion is created at $T \sim F_a$. Since then on the classical field $\langle a \rangle$ oscillate. Harmonic oscillator example suggests, $m_a^2 F_a^2$ is the energy density $\sim m_a \cdot$ (number density) which behaves like CDM.

Ref. [4] studies the axion energy density carefully with the axion field evolution equation for a time-varying Lagrangian and the adiabatic condition for the adiabatic invariant quantity, finding an overshoot factor of 1.8.

Thus, if axion is a significant CDM component of the universe [44], then it can be detected [45]. The experimental efforts for this is reviewed in [9]. Here, the feeble coupling is compensated by a huge number of axions, and there is a hope to detect a 10^{-5} eV axion. The limits from cosmic and astrophysical axion searches are shown in the previous figure, Fig. 3. The future ADMX and CARRACK will cover the interesting region.

CONCLUSION

I reviewed axions and the related issues on,

1. Solutions of the strong CP problem : the $m_u = 0$ possibility is ruled out now, the Nelson-Barr type still viable but without a compelling model yet, and the axion solution is most attractive and is not ruled out or may be very difficult to rule out.
2. If axions are discovered by cavity experiments, it will be the case of confirming instanton physics of QCD by experiments, which is most exciting.
3. Cosmology and astrophysics give bounds on the axion parameters. Maybe, axions are coming out from the WD cooling process. It is the first hint, in the middle of the axion window. A specific variant very light axion model has been constructed for $N_{DW} = \frac{1}{2}$.
4. With SUSY extension, O(GeV) axino can be CDM axino or decaying-to-CDM axino [43]. This kind of axino can produce the needed number of nonthermal neutralinos. In any case, to understand the strong CP with axions in the SUSY framework, the axino must be considered in the CDM discussion, which is presented at this conference [46].

ACKNOWLEDGMENTS

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